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Can Nearby Resonances Enhance $D^0 - \bar{D}^0$ Mixing?

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Abstract

We study the contributions of resonances to $D^0 - \bar{D}^0$ mixing. Both $Q\bar{Q}$ and hybrid $Q\bar{Q}G$ states are considered. Assuming reasonable values for the resonance parameters, we find relatively sizeable individual contributions to both Δm_D and $\Delta \Gamma_D$. We derive a variant of the GIM cancellation mechanism for the resonance amplitudes and show that broken $SU(3)$ can allow for appreciable residual effects. Additional input from meson spectroscopy and lattice gauge simulations will be needed to improve the accuracy of these predictions.

I. BASIC APPROACH

In this paper, we explore the possible enhancement in the mixing rate of neutral D mesons due to nearby resonances. It is interesting that the resonance mechanism, if operative, can be available only to D mesons. The light kaon lies below the resonance region and the heavy $B_{d,s}$ mesons lie above it. The dynamical mechanism of resonant enhancement constitutes an explicit violation of the quark-hadron duality assumption and could influence power counting rules built into the HQET estimate of Δm_D [1,2], which assumes a large energy gap between m_c and the scale Λ_{QCD} at which hadron dynamics is active. In addition, if a resonance is viewed as a single-particle intermediate state, then its contribution will be favored over that of multibody intermediate states by the $1/N_c$ counting rules.

For the remainder of this section, we continue the discussion of resonance contributions to mixing amplitudes. Section II concerns applications to $D^0 - \bar{D}^0$ mixing, Section III addresses the issue of $SU(3)$ multiplet structure and conclusions are presented in Section IV.

From standard perturbation theory, the ij^{th} element of the $D^0 - \bar{D}^0$ mass matrix can be represented as

$$\left[M - i\frac{\Gamma}{2} \right]_{ij} = \frac{1}{2m_D} \langle D_i^0 | \mathcal{H}_W^{\Delta C=2} | D_j^0 \rangle + \frac{1}{2m_D} \sum_I \frac{\langle D_i^0 | \mathcal{H}_W^{\Delta C=1} | I \rangle \langle I | \mathcal{H}_W^{\Delta C=1\dagger} | D_j^0 \rangle}{m_D^2 - m_I^2 + i\epsilon} . \quad (1)$$

The first term in the mass matrix expansion of Eq. (1) corresponds to the contribution of local $\Delta C = 2$ box and dipenguin operators. These are small in the Standard Model [3–5]. Next come the bilocal contributions which are induced by the insertion of two $\Delta C = 1$ operators. This class of terms might be enhanced by various nonperturbative effects, and therefore is of considerable interest. As follows from Eq. (1), one introduces a sum over all possible n -particle intermediate states allowed by the corresponding quantum numbers. For these continuum contributions, the summation in the second term of Eq. (1) takes the form of an integral over the energy variable. There will be a unitarity cut in the complex energy plane lying along the real axis and beginning at the two-pion threshold. The contribution from charged pseudoscalar two-body intermediate states was originally considered in Refs. [4,6] and estimated to be potentially large. However, it remains very difficult to reliably determine the total effect associated with $n \geq 2$ intermediate states due to the many decay modes present, each having unknown final state interaction (FSI) phases. For a recent attempt in this direction, see Ref. [7].

There are also the ‘single-particle’ effects arising both from bound states and from resonance intermediate states. A bound state contribution occurs as a pole on the real- E axis. A resonance contribution lies in the continuum and corresponds to a pole on an unphysical Riemann sheet. Its contribution will be a sharply peaked lorentzian profile of the discontinuity across the unitarity cut, much like that of a bound state. However, a resonance will contribute both to Δm_D and $\Delta \Gamma_D$. In principle, single-particle effects are rather simpler to analyze. The number of such intermediate states is constrained, and they can generally be estimated, at least roughly. As mentioned above the $1/N_c$ -counting rules, shown already to work reasonably well for the estimates of D -meson decay widths, favor a set of single-particle intermediate states, *i.e.* pole diagrams (*cf* Fig. 1) for self-energies of D_L . The light-meson and B -flavored meson single-particle contributions to $D^0 - \bar{D}^0$ mixing have already been analyzed [8]. In this paper we study the resonance sector.

Consider the special role of resonances. To begin, we express the collection of such contributions to Δm_D (upon neglecting CP-violation) as

$$\Delta m_D \Big|_{\text{tot}}^{\text{res}} = \frac{1}{2m_D} \sum_R \text{Re} \frac{\langle D_L | \mathcal{H}_W | R \rangle \langle R | \mathcal{H}_W^\dagger | D_L \rangle}{m_D^2 - m_R^2 + i\Gamma_R m_D} - (D_L \rightarrow D_S) . \quad (2)$$

The pseudoscalar 0^{-+} (scalar 0^{++}) intermediate states have $CP = -1$ ($CP = +1$) and contribute to the D_L (D_S) part of the above. If the mass of the resonance is not too far from the D -meson mass, an interesting effect occurs. To highlight the dependence on the resonance mass, temporarily consider just the energy denominator in Eq. (2). The contribution to the energy denominator of a light bound state (*e.g.* pions or kaons) of mass m is $P(m^2) \equiv 1/(m_D^2 - m^2) = m_D^{-2} + \mathcal{O}(m^2/m_D^2)$, which amounts to a suppression factor of order m_D^{-2} . By contrast, the energy denominator (*cf* Eq. (2)) for a resonance of mass m_R and width Γ_R will yield

$$\Delta m_D \Big|_R^{\text{res}} \propto \frac{m_D^2 - m_R^2}{(m_D^2 - m_R^2)^2 + \Gamma_R^2 m_D^2} , \quad \Delta \Gamma_D \Big|_R^{\text{res}} \propto -\frac{\Gamma_R m_D}{(m_D^2 - m_R^2)^2 + \Gamma_R^2 m_D^2} . \quad (3)$$

In the limit of a narrow resonance width, the expression for $\Delta \Gamma_D$ becomes proportional to the delta function $\delta(m_R - m_D)$, as expected. The effect of a finite width is to allow the resonance

to contribute at values $m_R \neq m_D$. The contribution to Δm_D vanishes at $m_R^2 = m_D^2$ since it undergoes a change of sign there. Considered as a function of the resonance mass m_R , the maximum effect occurs for $m_R^2 = m_D^2 \pm \Gamma_R m_D$ at which $\mathcal{R}e P(m_D^2 \pm \Gamma_R m_D) = \mp 1/(2\Gamma_R m_D)$. On the other hand, the resonance contribution to $\Delta\Gamma_D$ is maximized at the different value $m_R^2 = m_D^2$, but is still substantial at the values $m_R^2 = m_D^2 \pm \Gamma_R m_D$ which maximize the contribution to Δm_D . In particular, for both Δm_D and $\Delta\Gamma_D$, the m_D^{-2} dependence which would appear for a very light resonance has been replaced by $\Gamma_R m_D$. Thus, for a resonance sufficiently near the D meson the possibility exists for an enhancement factor of order $m_D/\Gamma_R \simeq 5 \rightarrow 15$ for both Δm_D and $\Delta\Gamma_D$ relative to an unenhanced pole contribution. Actually it even makes sense to broaden the term ‘nearby resonance’ to include $m_R > 1$ GeV since the m_D^{-2} suppression mentioned above will be largely overcome.

Also present in Eq. (2) are the D -to- R transition amplitudes. To obtain a quantitative description, we shall adopt as our $\Delta C = 1$ hamiltonian the phenomenological Bauer-Stech-Wirbel effective operator [10],

$$\mathcal{H}_W^{\text{BSW}} = \frac{G_F a_2}{\sqrt{2}} \bar{u}_k \Gamma_L^\mu c_k \left[V_{cd}^* V_{us} \bar{d}_j \Gamma_\mu^L s_j + V_{cs}^* V_{ud} \bar{s}_j \Gamma_\mu^L d_j + V_{cd}^* V_{ud} \bar{d}_j \Gamma_\mu^L d_j + V_{cs}^* V_{us} \bar{s}_j \Gamma_\mu^L s_j \right] , \quad (4)$$

where the constant $a_2 = -0.55 \pm 0.1$ is fixed from fitting nonleptonic D decays and it is understood that the operator of Eq. (4) is to be evaluated in vacuum saturation. In vacuum saturation, the contribution of resonance R to mixing will be proportional to the squared decay constant f_R^2 . This has two important consequences:

1. $Q\bar{Q}$ resonances having $J^P = 0^+$ will not contribute, as they occur in P-waves and thus have vanishing wave function at the origin. Although they could well be nonzero in a more general setting, their absence here suggests they would be suppressed.
2. We interpret $J^P = 0^-$ $Q\bar{Q}$ resonances with masses nearest the D as second radial excitations. We include also *first* radial excitations in our study due to their larger decay constants but omit radial excitations above the second.

II. CONTRIBUTIONS OF INDIVIDUAL RESONANCES

The fact that the most recent Particle Data Group compilation [9] cites many nonstrange and strange resonances in the mass region up to 2100 MeV strongly supports the premise of a resonance mechanism. Since experimental data are still relatively sparse in the mass region of the D system, however, additional spectroscopic knowledge of this energy range is needed.

In the following, we shall consider the effects of individual $Q\bar{Q}G$ or $Q\bar{Q}$ composites. Although individual contributions like these will, at least to some extent, be subject to GIM cancellations as other states are added in, such contributions nonetheless serve as useful indicators of what mixing signal to be reasonably expected. We employ Eq. (4) for the $Q\bar{Q}$ examples, while employing a largely phenomenological method for the $Q\bar{Q}G$ case.

$Q\bar{Q}$ Resonance

The mixing amplitudes induced by resonance R are

$$\Delta m_D^{(R)} = -C f_R^2 \frac{\mu_R(1 - \mu_R)}{(1 - \mu_R)^2 + \gamma_R^2}, \quad \Delta \Gamma_D^{\text{res}} = -C f_R^2 \frac{\mu_R \gamma_R}{(1 - \mu_R)^2 + \gamma_R^2}, \quad (5)$$

where $C \equiv 2m_D(G_F a_2 f_D \xi_d / \sqrt{2})^2$, the dimensionless quantities $\mu_R \equiv m_R^2/m_D^2$ and $\gamma_R \equiv \Gamma_R/m_D$ are the reduced squared-mass and width of the resonance, and from the unitarity of CKM matrix and neglecting ξ_b we have used $\xi_s \simeq -\xi_d$ where $\xi_i \equiv V_{ci}^* V_{ui}$ ($i = d, s$).

Since decay widths $\Gamma \simeq 0.20$ GeV are characteristic of resonances in the $1 \rightarrow 2$ GeV mass range, one expects that $1 \gg \gamma$ in applications of Eq. (5). We note that the two mass values $m_{\text{res}}(\text{GeV}) = 1.772, 1.973$ maximize the resonance contribution to $D^0 - \bar{D}^0$ mixing for fixed decay width and decay constant values. As to the dependence in Eq. (5) on the resonance decay constant one recalls $f_\pi \simeq 0.13$ GeV, $f_K \simeq 0.16$ GeV for the noncharm ground state and $f_D \simeq 0.2$ GeV, $f_{D_s} \simeq 0.28$ GeV [9,12] for the charm ground state. Excitations of the constituent quarks will reduce the wave function at the origin and thus decrease the corresponding decay constant. For $Q\bar{Q}$ radial excitations, we use hydrogen atom wave functions to provide a rough guide in estimating default values of f_R . For first radial excitations, we estimate $f_R \simeq 0.025$ GeV whereas for second radial excitations, we use $f_R \simeq 0.01$ GeV in our numerical work. We expect contributions from even higher radial excitations to be negligible. In our numerical work results are scaled with the square of the associated decay constant to allow for any future departures in assumed values or quantum number assignments of individual states.

We display in the Table various resonance contributions (assuming the $Q\bar{Q}$ description) to Δm_D and to $\Delta \Gamma_D$. The mass and decay width values are taken from the PDG listing. We do not intend our listing to be complete, but instead to indicate the magnitudes associated with contributions of this type. Although smaller than the current experimental limit [9,11] $|\Delta m_D|^{\text{expt}} < 1.3 \cdot 10^{-14}$ GeV, the values are larger than the contribution $\Delta m_D|^{\text{gnd state}} \simeq 3 \cdot 10^{-17}$ GeV from the set of pseudoscalar ground state mesons (π , K , η , η'). They also tend to dominate the short distance ‘box’ contributions $|\Delta m_D|^{\text{box}} \simeq 1.9 \times 10^{-17}$ GeV and $|\Delta \Gamma_D|^{\text{box}} \simeq 0.75 \times 10^{-17}$ GeV, where we have taken $m_c = 1.3$ GeV, $m_s = 0.2$ GeV, and refer effects of QCD radiative corrections to Ref. [2].

$Q\bar{Q}G$ Resonance

In vacuum saturation, contributions to Δm_D from $Q\bar{Q}$ intermediate states arise mainly from annihilation amplitudes. It is well known that such amplitudes are subject to helicity suppression. This is the same effect which influences the observed patterns of leptonic pion and kaon decay modes. It was noted long ago that the effects of helicity suppression can be (partially) lifted by soft-gluon emission. While this mechanism can be readily applied to inclusive heavy meson decay, it is difficult to see how to implement it to the mixing matrix elements of Eq. (2) if the intermediate states are of the $Q\bar{Q}$ variety. However, this difficulty is avoided if the intermediate state meson is a $Q\bar{Q}G$ hybrid state which involves a constituent gluon. The argument can be further extended to include penguin operators, thus introducing a long-distance counterpart of the dipenguin operator contribution [5,13]. It is plausible to assume that the gluon produced by a penguin (or any other operator) can form a quasibound state along with the $u\bar{u}$ quark-antiquark pair, thus producing a hybrid resonance (*cf* Fig. 2). These one-particle intermediate state contributions to Δm_D might be of importance because of the proximity of the anticipated [14,15] hybrid-meson mass with that of the D -meson.

In this regard, a particularly interesting candidate is the $\pi_H(1800)$ which has $J^{PC} = 0^{-+}$. On the basis of reports from several experimental groups [16] along with various quark model analyses [17], it is tempting to assign this particle as a hybrid. One can then estimate the contribution of $\pi_H(1800)$ to Eq. (2) provided that the mixing amplitude $g \equiv \langle D_L | H_w | \pi_H(1800) \rangle$ is known. This amplitude can be inferred from quark models or even better, determined phenomenologically by using available data on D decay rates. The idea is to search for common decay channels of D and $\pi_H(1800)$ where the $\pi_H(1800)$ contribution is manifest and then estimate the mixing amplitude from this. The situation is as depicted in Fig. 3, in which $D - \pi_H(1800)$ mixing is followed by $\pi_H(1800)$ decay.

It was noted in theoretical calculations [17] and hinted at experimentally that the decay rates $\pi_H(1800) \rightarrow \pi f_0(980)$, $\pi f_0(1300)$ are large for a hybrid $\pi_H(1800)$. Thus one can put an upper bound on the mixing amplitude g by introducing a model for the resonant decay of D -meson via $\pi_H(1800)$ [18],

$$\mathcal{M}_{D \rightarrow \pi f_0(980)} = \frac{g}{m_D^2 - m_{\pi_H}^2 + i\Gamma_{\pi_H} m_D} \mathcal{M}_{\pi_H \rightarrow \pi f_0(980)} \quad . \quad (6)$$

The partial decay width for $\pi_H \rightarrow \pi f_0(980)$ can be written as

$$\Gamma_{\pi_H \rightarrow \pi f_0(980)} = \frac{1}{16\pi m_{\pi_H}} \left| \mathcal{M}_{\pi_H \rightarrow \pi f_0(980)} \right|^2 \lambda_\pi \quad , \quad (7)$$

where λ_I for $I \rightarrow f_1 f_2$ is defined as

$$\lambda_I^2 \equiv \left[1 - \frac{(m_{f_1} + m_{f_2})^2}{m_I^2} \right] \left[1 - \frac{(m_{f_1} - m_{f_2})^2}{m_I^2} \right] \quad . \quad (8)$$

A similar formula exists for the $D \rightarrow \pi f_0(980)$ transition. In the simplest case, the total decay width of π_H can be saturated by the single partial decay width $\Gamma_{\pi_H \rightarrow \pi f_0(980)}$. This is a reasonable approximation as this decay mode becomes dominant for the hybrid π_H . Thus, using Eq. (7) along with an expression for $\Gamma_{D \rightarrow \pi f_0(980)}$, the mixing amplitude g can be estimated from

$$|g|^2 = \frac{m_D \Gamma_D \lambda_\pi}{m_{\pi_H} \Gamma_{\pi_H} \lambda_D} \cdot \left[(m_D^2 - m_{\pi_H}^2)^2 + \Gamma_{\pi_H}^2 m_D^2 \right] \quad . \quad (9)$$

Computing the mixing amplitude g using experimental data on the decay rate $\Gamma_{D \rightarrow \pi f_0(980)}$ and inserting it into Eq. (2), we estimate $|\Delta m_D|^{\pi_H(1800)} \leq 0.3 \times 10^{-16}$ GeV, comparable to the short distance result.

III. EFFECT OF MULTIPLY STRUCTURE

Resonances contributing as intermediate states to $D^0 - \bar{D}^0$ mixing are expected to occur as SU(3) flavor multiplets. The contribution of an entire multiplet will vanish in the limit of degenerate light-quark masses due to GIM cancellation. It is not clear how powerful the GIM suppression will be, as SU(3) is known to be badly broken in at least some D decays. For example, there is the experimentally measured ratio [9,19] $\Gamma_{D^0 \rightarrow K^+ K^-} / \Gamma_{D^0 \rightarrow \pi^+ \pi^-} \simeq 3$, which is unity in the $SU(3)$ limit. Theoretically, it has been suggested that such large

breaking is an accumulation of a number relatively minor effects whose ultimate impact is substantial. [20] At any rate, large SU(3) breaking could produce a loophole for evading GIM suppression.

Consider an octet of excited mesons π_H , K_H , \bar{K}_H , η_H . We use the subscript ‘H’ to represent *heavy* mesons, and denote the individual members of a resonance octet with the above flavor labels. We anticipate the presence of a ninth heavy meson η'_H to allow for mixing occurring with η_H . In principle, the mixing angle θ_H can be inferred from either from mass determinations or from two-photon branching ratios [9,21].

We write for the contribution of a mixed octet of resonances to Δm_D and Γ_D ,

$$\Delta m_D|_{\text{octet}}^{\text{res}} = \Delta m_D^{(K_H)} - \frac{1}{4}\Delta m_D^{(\pi_H)} - \frac{3\cos^2\theta_H}{4}\Delta m_D^{(\eta_H)} - \frac{1\sin^2\theta_H}{4}\Delta m_D^{(\eta'_H)} \quad (10)$$

$$\Delta \Gamma_D|_{\text{octet}}^{\text{res}} = \Delta \Gamma_D^{(K_H)} - \frac{1}{4}\Delta \Gamma_D^{(\pi_H)} - \frac{3\cos^2\theta_H}{4}\Delta \Gamma_D^{(\eta_H)} - \frac{1\sin^2\theta_H}{4}\Delta \Gamma_D^{(\eta'_H)} \quad , \quad (11)$$

where $\Delta m_D^{(i)}$ and $\Delta \Gamma_D^{(i)}$ are as in Eq. (5). The effect of SU(3) breaking can further be studied, say for Δm_D , by expressing the octet decay constant, mass and decay-width factors in Eq. (10) as

$$f_i = f_0 + \delta f_i \quad , \quad \mu_i = \mu_0 + \delta \mu_i \quad , \quad \gamma_i = \gamma_0 + \delta \gamma_i \quad (i = 1, \dots, 8) \quad , \quad (12)$$

where f_0, μ_0, γ_0 and $\delta f_i, \delta \mu_i, \delta \gamma_i$ represent respectively the SU(3)-invariant and SU(3)-breaking components. This allows for the possibility that the result will be generally influenced by SU(3)-breaking in the decay constant, mass and decay-width sectors. An expression valid to first order in symmetry breaking is

$$\Delta m_D \Big|_{\text{SU(3) brk}} = -\frac{C}{4} \left[F_f(f_0, \mu_0, \gamma_0) (4\delta f_{K_H} - \delta f_{\pi_H} - 3\delta f_{\eta_H}) \right. \\ \left. + F_\mu(f_0, \mu_0, \gamma_0) (4\delta \mu_{K_H} - \delta \mu_{\pi_H} - 3\delta \mu_{\eta_H}) + F_\gamma(f_0, \mu_0, \gamma_0) (4\delta \gamma_{K_H} - \delta \gamma_{\pi_H} - 3\delta \gamma_{\eta_H}) \right] \quad , \quad (13)$$

where for simplicity η_H - η'_H mixing is ignored and we define

$$F_f \equiv f_0 \frac{2\mu_0(1-\mu_0)}{(1-\mu_0)^2 + \gamma_0^2} \quad , \quad F_\mu \equiv f_0^2 \frac{(1-\mu_0)^2 + (1-2\mu_0)\gamma_0^2}{[(1-\mu_0)^2 + \gamma_0^2]^2} \quad , \quad F_\gamma \equiv -f_0^2 \frac{2\gamma_0\mu_0(1-\mu_0)}{[(1-\mu_0)^2 + \gamma_0^2]^2} \quad . \quad (14)$$

Due to the relative lack of data, it is not possible at this time to provide a unique analysis of the above relations. Either a small or large effect could emerge, for example:

1. In the absence of η_H - η'_H mixing, the combination $4\delta \mu_{K_H} - \delta \mu_{\pi_H} - 3\delta \mu_{\eta_H}$ vanishes by virtue of the Gell Mann-Okubo formula and the remaining dependence in Eq. (13) vanishes with the choice $\mu_0 = 1$ (implying $F_f = F_\gamma = 0$).
2. The choice $\mu_0 = 1 - \gamma_0$ (with $1 \gg \gamma_0$) yields

$$F_f \simeq \frac{f_0}{\gamma_0} \quad , \quad F_\mu \simeq \frac{f_0^2}{2\gamma_0} \quad , \quad F_\gamma \simeq -\frac{f_0^2}{2\gamma_0} \quad .$$

Here the net effect appears in terms of fractional changes in decay constants and decay rates,

$$|\Delta m_D|^{\text{SU}(3) \text{ brk}} \simeq 0.2 \times 10^{-14} \frac{f_0^2}{\gamma_0} \text{ GeV} \left| \frac{4\delta f_{K_H} - \delta f_{\pi_H} - 3\delta f_{\eta_H}}{f_0} - \frac{4\delta \gamma_{K_H} - \delta \gamma_{\pi_H} - 3\delta \gamma_{\eta_H}}{2\gamma_0} \right|.$$

In the first of the above items, the vanishing of $\Delta m_D|^{\text{SU}(3) \text{ brk}}$ to first order in SU(3) symmetry breaking occurs for a special parameter choice and is clearly more an exception than a rule. The second item has the SU(3) degenerate mass set at a lower value (still with no η_H - η'_H mixing) and a nonzero effect will generally occur, although to be more quantitative would require additional experimental input. The presence of large symmetry breaking might necessitate a treatment beyond the first-order relations given above.

IV. CONCLUDING REMARKS

The motivation most often cited in searches for $D^0 - \bar{D}^0$ mixing lies with the possibility of observing a signal from new physics which dominates that from the Standard Model. The best experimental limit, recently obtained by E791 [11], is well beneath existing estimates of the Standard Model value. There are plans to improve on the E791 determination, both at B-factories [22] and at hadron colliders [23]. In addition, preliminary plans at Jefferson Lab to build a new experimental hall and simultaneously to raise the beam energy suggest the possibility for $D^0 - \bar{D}^0$ mixing studies at that facility as well. For all such efforts, it will be crucial to understand the magnitude of the mixing amplitude from the Standard Model.

In this paper, we have studied the set of potentially significant contributions to $D^0 - \bar{D}^0$ mixing from pseudoscalar resonances. We have shown how an enhancement of order m_D/Γ_R can arise from a resonance whose mass lies within several decay widths of m_D and have also have pointed out the importance of lighter resonances due to decay constant dependence in the mixing amplitude. In order to obtain a more detailed understanding of the resonance scenario, we have considered contributions from both traditional $\bar{Q}Q$ resonances as well as an exotic $\bar{Q}QG$ hybrid. Effects of order 10^{-16} GeV are possible for both Δm_D and $\Delta \Gamma_D$. In addition, it would appear possible or even likely in the resonance mechanism that $|\Delta \Gamma_D/\Delta m_D| \simeq 1$ or even larger. This calls into question the usual assumption, that $\Delta m_D \gg \Delta \Gamma_D$, made in searches for CP violation in D-decay using time-dependent measurements [24]. We shall consider generalizations of our approach and implications of our findings *vis-a-vis* CP-violating signals in a separate publication.

Of course, efforts such as this are severely hampered by a lack of knowledge regarding the properties of mesons lying in the $1.6 \rightarrow 2.1$ GeV mass range. Two kinds of additional input would be of significant value to this subject. From experimentalists could come a more complete listing of resonance mass and decay-width parameters. Information about various decay modes could allow a distinction between the $\bar{Q}Q$ and $\bar{Q}QG$ descriptions. The lattice-gauge community could supply numerical estimates of both decay constants of excited mesons and also, as a test of vacuum saturation, matrix elements like $\langle a_0 | H_W | D^0 \rangle$. As such valuable information becomes available, we can anticipate real progress in this area.

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FIGURES

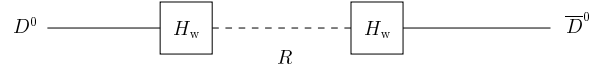


FIG. 1. Contribution of resonance R to the D^0 -to- \bar{D}^0 matrix element.

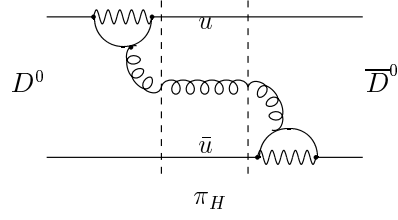


FIG. 2. $Q\bar{Q}G$ intermediate state.

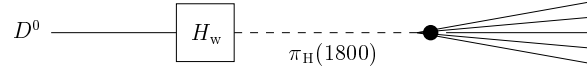


FIG. 3. Hybrid contribution to D^0 decay.

TABLES

Table: Magnitudes of Pseudoscalar Resonance Contributions.		
Resonance	$ \Delta m_D \times 10^{-16} \text{ (GeV)}$	$ \Delta \Gamma_D \times 10^{-16} \text{ (GeV)}$
$K(1460)$	$\sim 1.24 (f_{K(1460)}/0.025)^2$	$\sim 0.88 (f_{K(1460)}/0.025)^2$
$\eta(1760)$	$(0.77 \pm 0.27) (f_{\eta(1760)}/0.01)^2$	$(0.43 \pm 0.53) (f_{\eta(1760)}/0.01)^2$
$\pi(1800)$	$(0.13 \pm 0.06) (f_{\pi(1800)}/0.01)^2$	$(0.41 \pm 0.11) (f_{\pi(1800)}/0.01)^2$
$K(1830)$	$\sim 0.29 (f_{K(1830)}/0.01)^2$	$\sim 1.86 (f_{K(1830)}/0.01)^2$